

Optical-infrared flares and radio afterglows from the tidal disruption of Jovian planets by their host star

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ABSTRACT

When a Jovian planet gets sufficiently close to its host star to be tidally disrupted, its debris stream deposits energy on the star’s surface, producing an expanding bubble of hot plasma. We study the radiation from the bubble and show that it includes optical-infrared prompt emission and a subsequent radio afterglow. The prompt emission from M31 and Large Magellanic Cloud is detectable by optical-near infrared transient surveys with a large field of view at an event rate of a few events per year. The subsequent radio afterglows are detectable for 10^{3-4} years.

Key words: radiation mechanisms: non-thermal — radiation mechanisms: thermal — planet-star interactions — infrared: planetary systems — radio continuum: planetary systems

1 INTRODUCTION

A substantial fraction of gaseous planetary-mass objects might be ingested by the central stars (Machida et al. 2010, 2011; Inutsuka 2012; Vorobyov & Basu 2010, 2015, and references therein). Realistic non-ideal magnetohydrodynamics simulations have shown that protoplanetary disks are initially massive enough to produce multiple Jupiter-mass planets via gravitational instability (Inutsuka et al. 2010; Machida et al. 2011; Tsukamoto et al. 2015). These massive planets may survive in the subsequent era, during which planets gravitationally interact or collide with each other to produce hot Jupiters and highly eccentric planets (e.g., Ida & Lin 2004; Chatterjee et al. 2008; Ford & Rasio 2008). A large fraction of the hot Jupiters which migrate to the vicinity of the central stars are either consumed (Sandquist et al. 1998) or tidally disrupted (Gu et al. 2003) by the host star. Stars without detected hot Jupiters might have already ingested them (Rice et al. 2008; Inutsuka 2009; Ogiwara et al. 2013, 2014). Present-day hot Jupiters could secularly enlarge their eccentricity to reach their host stars by a process like the Kozai-Lidov mechanism (Kozai 1962; Lidov 1962). Several ways for detecting stellar ingestion of planets have been proposed (e.g., Sandquist et al. 1998; Jackson et al. 2009; Teitler & Konigl 2014; Matsakos & Konigl 2015). If a star ingests planets on average N_i times during its life, the total event rate in the

Galaxy is estimated to be $\text{SFR} \times N_i / \langle m \rangle \sim 5N_i \text{ yr}^{-1}$, where $\text{SFR} \approx 1M_\odot \text{ yr}^{-1}$ is the star formation rate of the Milky Way (Robitaille & Whitney 2010) and $\langle m \rangle \approx 0.2M_\odot$ is the average stellar mass (Kroupa 2001; Chabrier 2003).

In this paper, we calculate the radiation expected at the moment of Jovian planet ingestion and the subsequent afterglow phase. When a planet is tidally disrupted, the resultant tidal stream hits the surface of its host star, releasing gravitational energy which is converted into an expanding bubble of hot plasma (§ 2). The expanding bubble generates an optical-infrared flare (§ 3). Subsequently, the material interacts with the circumstellar matter (CSM), generating a shock that accelerates electrons to relativistic energies, which in turn produce synchrotron radiation (§ 4).

So far, tidal disruption events (TDEs) have been widely investigated for the case of the supermassive black holes (SMBHs) hosting stars, providing a distinct opportunity to probe dormant SMBHs in inactive galaxies. Most TDEs take place when a distant star is perturbed into a parabolic orbit approaching close enough to the SMBH to be ripped apart by its tidal force. The subsequent accretion of stellar debris falling back to the SMBH produces a characteristic flare, which decays with the $-5/3$ power law of time (Rees 1988; Phinney 1989), with a luminosity large enough to exceed the Eddington luminosity for a timescale of weeks (Evans & Kochanek 1989; Guillochon et al. 2011). Candidates for TDEs have also been observed in various wavebands (see Komossa 2015, for a review). Such events can also occur in a multiple exoplanet system by the same mechanism

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as the SMBH context except for that a planet is tidally disrupted by a central star.

2 INITIALLY RELEASED ENERGY

Most extrasolar planets reside in eccentric orbits (Marcy et al. 2005; Winn & Fabrycky 2015). Some planets are subject to rapid orbital change or migration by the interaction with a gaseous disk (Rice et al. 2008), planet-planet scattering (Rasio & Ford 1996; Weidenschilling & Marzari 1996), or the Kozai-Lidov mechanism (Ford et al. 2000; Fabrycky & Tremaine 2007; Li et al. 2015).

Following these dynamical processes, the pericentre distance R_p finally becomes shorter than the tidal disruption radius of a planet: $R_T = (M_*/m_{pl})^{1/3} r_{pl}$, where M_* is a mass of the star hosting the planet with a mass m_{pl} and a radius r_{pl} . Then, the planet is tidally disrupted (e.g., Faber et al. 2005). Subsequently, the planet debris falls back and is expected to hit the stellar surface because $R_* \lesssim R_p \lesssim R_T$, where R_* is the stellar radius. Note that $R_T \sim R_*$ for Sun-like G-type stars and Jovian mass planets, whereas $R_T > R_*$ for K- and M-stars (e.g., Rappaport et al. 2013).

After the tidal disruption of a planet on a parabolic orbit, the debris mass is distributed around zero specific energy (see Figure 3 of Evans & Kochanek 1989), characteristic of a parabolic orbit. If the planet originally approaches the star on an eccentric orbit, however, the mass distribution of the disrupted planet would shift to negative specific energy. This provides the condition that an eccentric, bound orbit would follow tidal disruption: $0 \leq e < e_{crit}$, where $e_{crit} = 1 - (2/\beta)(m_{pl}/M_*)^{1/3}$ with a penetration factor $\beta = R_T/R_p$ (Hayasaki et al. 2013). In the eccentric TDEs, the mass fallback rate is significantly deviated from the scaling $t^{-5/3}$ as expected for the parabolic TDEs, because the mass falls back to the star over a much shorter accretion time $\Delta t \sim m_{pl}/\dot{M}$ (Hayasaki et al. 2013, 2015; Bonnerot et al. 2015).

Half of the gravitational binding energy of the accreting matter is released as the thermal energy at a rate, $L_{th} = GM_*\dot{M}/2R_*$. Its magnitude for eccentric TDEs is therefore much larger than the parabolic TDE case. However, the total injected thermal energy, $E_{th,i} \sim L_{th}\Delta t$, eventually takes the same value between eccentric and parabolic TDEs:

$$E_{th,i} \sim \frac{GM_*m_{pl}}{2R_*} = 1.8 \times 10^{45} \xi_* \left(\frac{m_{pl}}{m_J} \right) \text{ erg} , \quad (1)$$

where $\xi_* = (M_*/M_\odot)/(R_*/R_\odot)$ and $m_J = 1.898 \times 10^{30} \text{ g}$ is the Jupiter mass. Note that ξ_* is of order unity for main sequence stars with $M_* \lesssim 2M_\odot$ (e.g., Torres et al. 2010; Eker et al. 2015).

3 PROMPT EMISSION FROM EXPANDING PLASMA BUBBLE

When the planetary debris hits the surface of its host star, a thermal energy $E_{th,i}$ is released into a volume of radius R_i over a short time. We approximate $R_i \sim r_{pl}$ and the number density of the confined gas, $n_i \sim 3m_{pl}/4\pi\mu_p R_i^3$, where μ_p is the proton mass. The plasma bubble is optically thick to its thermal photons since initial optical depth is estimated

as $\tau_i = n_i\sigma_T R_i = 3.7 \times 10^9 (m_{pl}/m_J)(R_i/r_J)^{-2}$, where σ_T is the Thomson cross section and $r_J = 7.0 \times 10^9 \text{ cm}$ is the radius of Jupiter. Subsequently, the bubble expands due to its thermal pressure. Here, we focus on a simple estimate of the luminosity of the resulting emission by the bubble. For simplicity, we assume that the gas has uniform density and temperature, and a homologous velocity profile.

The bubble is matter dominated, so that the initial temperature is given by

$$T_i \sim \frac{GM_*\mu_p}{3k_B R_*} = 7.7 \times 10^6 \xi_* \text{ K} , \quad (2)$$

where k_B is Boltzmann's constant. The initial radiation energy, $E_{rad,i} \sim (aT_i^4)(4\pi R_i^3/3) \sim 3.8 \times 10^{43} \xi_*^4 (R_i/r_J)^3 \text{ erg}$, where a is the radiation energy constant, is much smaller than $E_{th,i}$ in Eq. (1). The material is initially opaque to its own thermal photons, allowing very little internal energy to escape from its surface. The hot plasma expands adiabatically with an expansion speed comparable to the escape velocity of the star, $v_{esc,*} = (2GM_*/R_*)^{1/2} = 6.2 \times 10^7 \xi_*^{1/2} \text{ cm s}^{-1}$. The temperature declines adiabatically as a function of radius R , as $T(R) = T_i(R/R_i)^{-2}$, while the gas number density is given by $n(R) = n_i(R/R_i)^{-3}$.

The expanding bubble becomes optically thin when it cools below the hydrogen recombination temperature of $\sim 10^4 \text{ K}$. The photosphere radius R_{ph} is determined by the condition that the photon diffusion time $t_{diff} = n_e\sigma_T R^2/c$, where $n_e = n_e(R)$ is the number density of free electrons, is equal to the expansion time $t_{exp} = R/v_{esc,*}$ at R_{ph} . We define $x = n_e/n(R)$ as the ionization degree at radius R . The Saha equation for hydrogen,

$$\frac{1-x}{x^2} = n(R) \left(\frac{h^2}{2\pi\mu_e kT(R)} \right)^{3/2} \exp \left(\frac{13.6 \text{ eV}}{kT(R)} \right) , \quad (3)$$

combined with $n/n_i = (T/T_i)^{3/2}$ and $t_{diff}/t_{exp} = x\tau_i(T/T_i)(v_{esc,*}/c) = 1$, can provide the equation for the temperature T_{ph} at the photosphere, to numerically find $T_{ph} \approx 7200 \text{ K}$. Due to the exponential term in Eq. (3), the value of T_{ph} is almost independent of the initial state of the bubble. The photosphere radius is therefore,

$$\begin{aligned} R_{ph} &= R_i \left(\frac{T_i}{T_{ph}} \right)^{1/2} \\ &= 2.3 \times 10^{11} \xi_*^{1/2} \left(\frac{R_i}{r_J} \right) \left(\frac{T_{ph}}{7200 \text{ K}} \right)^{-1/2} \text{ cm} , \end{aligned} \quad (4)$$

and the ionization degree at $R = R_{ph}$ is $x(R_{ph}) = 1.4 \times 10^{-4} (m_{pl}/m_J)^{-1} (R_i/r_J)^2 \xi_*^{-1/2} (T_{ph}/7200 \text{ K})^{-1}$. The observer would detect blackbody radiation with a temperature T_{ph} and a peak bolometric luminosity,

$$\begin{aligned} L_p &= 4\pi R_{ph}^2 \sigma T_{ph}^4 \\ &= 1.0 \times 10^{35} \xi_* \left(\frac{R_i}{r_J} \right)^2 \left(\frac{T_{ph}}{7200 \text{ K}} \right)^3 \text{ erg s}^{-1} , \end{aligned} \quad (5)$$

where σ is the Stefan-Boltzmann constant. The peak flux density at frequency $\nu = \nu_{14} \times 10^{14} \text{ Hz}$ and a distance $d = 1 d_{kpc}$ kpc from the source is then

$$\begin{aligned} F_\nu^p &= \frac{L_p}{4\pi d^2} \frac{15}{\pi^4 \nu} \left(\frac{h\nu}{kT_{ph}} \right)^3 f \left(\frac{h\nu}{kT_{ph}} \right) \\ &= 38 \xi_* \left(\frac{R_i}{r_J} \right)^2 \nu_{14}^2 d_{kpc}^{-2} f \left(\frac{h\nu}{kT_{ph}} \right) \text{ mJy} , \end{aligned} \quad (6)$$

Table 1. Predicted optical/infrared peak flux density of the prompt emission. The unabsorbed observed peak flux, F_ν^p , is for the distance $d = 10$ kpc from the source with $\xi_* = 1$, $R_i = r_J$, and $T_{ph} = 7200$ K.

Filter	λ [nm]	ν [10^{14} Hz]	F_ν^p [mJy]
g'	475	6.3	0.97
r'	622	4.8	1.2
i	763	3.9	1.2
y	1020	2.9	1.1
J	1220	2.5	0.91
H	1630	1.8	0.66
K	2190	1.4	0.44
L	3450	0.87	0.21
M	4750	0.63	0.12
N	10500	0.29	0.028

where $f(y) = y(e^y - 1)^{-1}$. Table 1 provides the flux density in various observation bands. Note that since $\xi_* \approx 1$ for $M_* \lesssim 2M_\odot$, the observed flux hardly depends on stellar properties.

The typical duration of the transient is comparable to the dynamical time-scale,

$$\Delta T \sim \frac{R_{ph}}{v_{esc,*}} = 3.7 \times 10^3 \left(\frac{R_i}{r_J} \right) \left(\frac{T_{ph}}{7200 \text{ K}} \right)^{-1/2} \text{ s}, \quad (7)$$

so that the total emission energy is

$$E_{rad} \sim L_p \Delta T = 3.7 \times 10^{38} \xi_* \left(\frac{R_i}{r_J} \right)^3 \left(\frac{T_{ph}}{7200 \text{ K}} \right)^{5/2} \text{ erg}, \quad (8)$$

which is much smaller than the initial internal energy of the bubble. Hence, almost all the initial energy transforms to kinetic energy and gets dissipated when the expanding material interacts with the CSM.

At the moment of tidal disruption, the planet is expected to be vertically collapsed (e.g., Kobayashi et al. 2004; Guillochon et al. 2009). The work done by the tidal force from the star is estimated to be $W \sim (GM_* m_{pl}/R_p^2)(r_{pl}/R_p)(r_{pl}/2) \sim \beta^3 G m_{pl}^2 / 2 r_{pl}$, where $\beta = R_T/R_p \sim 1$ is a penetration factor. If the collapsed matter is thermalized and half of this energy is released, then the initial thermal energy becomes $E_{th,i} \sim 8.6 \times 10^{42} \beta^3 (m_{pl}/m_J)^{5/3} \text{ erg}$, where we use an approximate relation, $(r_{pl}/r_J) \approx (m_{pl}/m_J)^{1/3}$ (Rappaport et al. 2013). This is about two orders of magnitude smaller than the energy considered in Eq. (1). The initial temperature is then estimated as $T_i \sim 3.7 \times 10^4 \beta^3 (m_{pl}/m_J)^{2/3} \text{ K}$, which is of order the recombination temperature. The bubble expands at a speed comparable to the free-fall velocity of the planet, $v_{ff,pl} \sim (G m_{pl}/r_{pl})^{1/2} = 4.3 \times 10^6 (m_{pl}/m_J)^{1/2} (r_{pl}/r_J)^{-1/2} \text{ cm s}^{-1}$. Similarly to the previous calculation, we derive a temperature of $T_{ph} \approx 8100 \text{ K}$ at the photosphere radius $R_{ph} = 1.5 \times 10^{10} (m_{pl}/m_J)^{2/3} \beta^{3/2} (T_{ph}/8100 \text{ K})^{-1/2}$. The peak bolometric luminosity is then $L_p \sim 6.7 \times 10^{32} \beta^3 (m_{pl}/m_J)^{4/3} (T_{ph}/8100 \text{ K})^3 \text{ erg s}^{-1}$. This is smaller than the main prompt emission (see Eq. 5), however, it could be larger for larger β and/or m_{pl} , in which case the emission could be detectable as a precursor arising before the prompt emission.

4 RADIO AFTERGLOW

The expanding plasma maintains a constant velocity $v_{esc,*}$ out to the deceleration radius,

$$R_{dec} = \left(\frac{3m_{pl}}{4\pi\mu_p n_c} \right)^{1/3} = 6.5 \times 10^{17} n_c^{-1/3} \left(\frac{m_{pl}}{m_J} \right)^{1/3} \text{ cm}, \quad (9)$$

where n_c is the density of the CSM. The bubble reaches this radius after $t_{dec} = R_{dec}/v_{esc,*} = 3.3 \times 10^2 \xi_*^{-1} n_c^{-1/3} (m_{pl}/m_J)^{1/3} \text{ yr}$. During the expansion, the flow interacts with the CSM, generating an external shock with a Mach number of ~ 60 . Electron acceleration at the shock results in radio synchrotron emission. The emission lasts until the shock velocity declines to $\sim 1 \times 10^7 \text{ cm s}^{-1}$, below which the ionization of the acceleration region drops rapidly (Shull & McKee 1979), so that wave damping due to collisions with neutral atoms prevents electrons from being accelerated at the shock front (Drury et al. 1996; Bykov et al. 2000). Assuming the Sedov solution after a time t_{dec} ($R \propto t^{2/5}$ and $v \propto t^{-3/5}$), we estimate the epoch t_{end} at which particle acceleration ceases to be $t_{end} \sim 6.9 \times 10^3 \xi_*^{-1/6} n_c^{-1/3} (m_{pl}/m_J)^{1/3} \text{ yr}$, corresponding to the radius $R_{end} \sim 2.2 \times 10^{18} \xi_*^{1/3} n_c^{-1/3} (m_{pl}/m_J)^{1/3} \text{ cm}$. When the acceleration stops, the high-energy electrons starts to escape from the shocked region with an escape time of $\sim 1\text{--}10 \text{ yr}$, resulting in rapid fading of the emission.

Next we provide a simple estimate of the observed flux and surface brightness of the radio synchrotron emission from the expanding bubble of radius R . The total number of nonthermal electrons, N_e , is a fraction η_e of the number of particles originating from the upstream region of the shock over a dynamical time $t_{dyn} = R/v$, where $v \sim v_{esc,*}$ is the expansion speed, that is, $N_e = \eta_e (4\pi R^2 n_c v) t_{dyn} = 4\pi \eta_e n_c R^3$. We assume a single power-law form of the nonthermal electron distribution, $N(\gamma) \propto \gamma^{-p}$, for $\gamma_m < \gamma < \gamma_M$, where γ is the electron Lorentz factor. If the dynamical time is sufficiently long ($t_{dyn} \gtrsim 2 \times 10^3 \text{ yr}$), the maximum Lorentz factor γ_M is determined by the balance of the acceleration time and the synchrotron cooling time (e.g., Yamazaki et al. 2004, 2006, 2015), yielding $\gamma_M \approx 1 \times 10^8 (\xi_*/f B_{-5})^{1/2} (v/v_{esc,*})$, where $B_{-5} = (B/10 \mu\text{G})$ is the post-shock magnetic field strength, and f is a numerical factor of order unity which is determined by the properties of scattering waves and shock geometry at the acceleration site. Thus, we find that radio-emitting electrons have a much smaller Lorentz factor than γ_M .

The synchrotron cooling time of radio emitting electrons is much longer than the dynamical time. Hence, the observed flux density at frequency ν is (e.g., Sari, Piran, & Narayan 1998)

$$\begin{aligned} F_\nu &\sim \frac{N_e \mu_e c^2 \sigma_T B}{4\pi d^2 3e} \left(\frac{\nu}{\nu_m} \right)^{(1-p)/2} \\ &= 4.0 \times 10^2 \frac{n_c \eta_{-5} B_{-5}}{d_{\text{kpc}}^2} \left(\frac{R}{10^{18} \text{ cm}} \right)^3 \left(\frac{\nu}{\nu_m} \right)^{(1-p)/2} \text{ Jy}, \end{aligned} \quad (10)$$

where $\eta_{-5} = (\eta_e/10^{-5})$ and $\nu_m = 28 B_{-5} \gamma_m^2 \text{ Hz}$ is the characteristic synchrotron frequency from electrons with a minimum Lorentz factor γ_m . Note that for our parameters, $(\nu/\nu_m)^{(1-p)/2}$ is much lower than unity (see Table 2). The

Table 2. Radio surface brightness at a radius $R = 10^{18}$ cm and frequency $\nu = 1$ GHz for power-law electron distribution with index p . Other parameters are taken as $n_c = \eta_{-5} = B_{-5} = \gamma_m = 1$.

p	$(\nu/\nu_m)^{(1-p)/2}$	$S_{\nu=1\text{GHz}}$ [Jy sr $^{-1}$]
2.0	1.7×10^{-4}	2.0×10^5
2.1	7.0×10^{-5}	8.4×10^4
2.2	2.9×10^{-5}	3.5×10^4
2.3	1.2×10^{-5}	1.5×10^4
2.4	5.2×10^{-6}	6.2×10^3
2.5	2.2×10^{-6}	2.6×10^3
3.0	2.8×10^{-8}	34

surface brightness, S_ν , is the flux density divided by the solid angle of the source, $\Omega \sim \pi(R/d)^2$, yielding

$$S_\nu \sim 1.2 \times 10^9 n_c \eta_{-5} B_{-5} \frac{R}{10^{18} \text{cm}} \left(\frac{\nu}{\nu_m} \right)^{(1-p)/2} \text{Jy sr}^{-1}. \quad (11)$$

Table 2 shows the results for different values of p with fixed parameters, $n_c = \eta_{-5} = B_{-5} = \gamma_m = 1$. For comparison, the surface brightness of the faintest Galactic supernova remnants is $\sim 10^4$ Jy sr $^{-1}$ (Arbutina & Urošević 2005), of the same order as the typical value of the diffuse Galactic radio emission (e.g., de Oliveira-Costa et al. 2008).

5 DISCUSSION

While the prompt emission flare from the tidal disruption of a planet cannot be detected in the optical band due to Galactic dust extinction, the unabsorbed flux in the K-band is ~ 0.2 mJy at the distance of 10 kpc, potentially detectable with infrared sky surveys, such as UKIRT Infrared Deep Sky Survey (UKIDSS: Hewett et al. 2006; Lawrence et al. 2007) and VISTA Variables in the Via Lactea (VVV: Minniti et al. 2010). Unfortunately, the expected event rate ~ 5 events yr $^{-1}$ ($N_i \approx 1$) for the entire Galactic plane (see § 1), is too small for detectability by current transient surveys. The duration ΔT of the expected prompt flares could be comparable to that of superflares of stars (e.g., Schaefer et al. 2000; Maehara et al. 2012). The total emission energy in the optical band is about an order of magnitude larger than the largest superflares (Shibayama et al. 2013; Balona 2015).

The Andromeda galaxy, M31, located at the distance of 0.78 Mpc (Stanek & Garnavich 1998), has a similar mass and star formation rate to the Milky Way (Williams 2003), and hence a similar planetary disruption event rate. Unlike the Milky Way case, telescopes with a large field of view can cover the entire volume of M31. The expected AB magnitudes are 25.9 and 25.7 mag at g and r bands, respectively, for our fiducial parameters. The fluxes become higher for planets more larger than Jupiter. Using the pixel lensing technique (i.e., differential image photometry: Crotts 1992; Baillon et al. 1993; Tomaney & Crotts 1996; Calchi Novati 2010), those transients could be detected by future instru-

ments with better sensitivity like Subaru Hyper Suprime-Cam¹.

Events in the Large Magellanic Cloud (LMC) could also be detected. For the distance of 48.5 kpc to LMC (Macri et al. 2006), the expected AB magnitudes are 19.9 and 19.6 mag at g and r bands respectively for our fiducial parameters, which are detectable by current surveys like PAndromeda (Lee et al. 2012). The event rate is only slightly smaller than M31 because the total star formation rate of the LMC is $\approx 0.4 M_\odot \text{yr}^{-1}$ (Harris & Zaritsky 2009).

If the electron index p is smaller than about 2.3, the radio afterglow would be detectable for $\sim 10^{3-4}$ yr after the disruption event. The radio surface brightness is expected to be lower than young supernova remnants because the magnetic field is weak. As a result, the surface brightness and diameter of the source would be distinguishable from the values expected for supernova remnants (Arbutina & Urošević 2005).

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¹ <http://www.naoj.org/Projects/HSC/index.html>

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